

B-Decay Signatures of Warped Top-Condensation

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We point out that the light radion ϕ in a recently proposed Warped Top-Condensation Model (WTCM), can provide distinct signatures in $b \rightarrow s\phi$, where the on-shell ϕ can decay with displaced vertices. We find that some of the parameter space of these models is constrained by B -meson and astrophysical data. Future B -decay measurements can lead to the discovery of the WTCM.

The mechanism responsible for electroweak symmetry breaking (EWSB) remains unknown. In the Standard Model (SM), the vacuum expectation value (vev) of the Higgs scalar results in EWSB. However, this minimal picture is unstable against quadratically divergent quantum corrections that would naturally pull the Higgs vev near a large UV cutoff scale. This puzzle is often referred to as the hierarchy problem, with the UV scale taken to be the Planck mass $\bar{M}_P \simeq 2 \times 10^{18}$ GeV.

In recent years, warped 5D models of the Randall-Sundrum (RS) type [1] have been developed as a possible framework to address the hierarchy. These models are based on a slice of 5D Anti de Sitter (AdS) space-time of curvature k , truncated by 4D boundaries, sometimes called the UV and the IR branes. The hierarchy is solved if the Higgs is localized near the IR brane and the distance L between the branes is stabilized such that $kL \sim 35$. The stabilization leads to the appearance of a weak scale scalar, the radion ϕ .

Even though warped models can address the hierarchy, another possibility that can achieve the same end is to remove fundamental scalars, and in particular the Higgs, from the theory. In this case, one can still have EWSB through the condensation of fermions, in analogy with chiral symmetry breaking in QCD. Such models generically require the appearance of strong interactions near the weak scale, as is the case in technicolor scenarios. However, the requisite strong interactions can also be provided by the exchange of Kaluza-Klein (KK) states in warped RS-type models. Even though these models are conjectured to be dual to 4D strongly-coupled theories [2, 3], based on the AdS/Conformal Field Theory (CFT) correspondence [4], the geometric 5D approach is interesting in its own right as a quantitative framework.

A recently proposed model employs strong interactions, mediated by KK gluons, between top quarks to provide effective dim-6 operators that lead to top condensation and EWSB [5] (for a related but different approach see Ref. [6]). In this model, the size L of the extra dimension is set by the condensation and is a prediction of the theory. We will refer to this framework as the Warped Top-Condensation Model (WTCM), in the fol-

lowing. However, this setup yields very heavy KK modes, of order 30 TeV, which are well-beyond the reach of the CERN Large Hadron Collider (LHC). In this model, accommodating the measured top mass requires the addition of extra TeV scale quarks, charged under SU(3) color, which can be searched for. However, this is the only direct LHC signature of this model, apart from a composite Higgs doublet of mass around 500 GeV. Nonetheless, there is another feature of this setup that can in principle yield distinct signatures, namely the appearance of a GeV-scale radion field ϕ [5]. Theories whose main new physics signature is a light dilaton have also been studied in [7], although in a different region of parameter space.

In this work, we point out that for $m_\phi \sim 1$ GeV, one can look for signatures of the WTCM in B -decays. In particular, decay modes of the type $b \rightarrow s\phi$ can be allowed on-shell and result in displaced vertex signatures, through ϕ -decays. We show that both the loop-level and the tree-level FCNC effects can lead to non-trivial constraints on the model and can also result in its potential discovery, using B -meson decay data. We will next briefly present the essential properties of the radion relevant to our framework.

The radion ϕ arises as the pseudo-Goldstone scalar associated with breaking of dilatation invariance, after the stabilization of the extra dimension. In 5D warped models [8], ϕ couples to the trace of energy momentum tensor and its coupling is suppressed by the scale

$$\Lambda_\phi = \sqrt{6M_5^3/k} e^{-kL}, \quad (1)$$

where M_5 is the 5D fundamental scale and $\bar{M}_P^2 = M_5^3/k$, up to tiny $\mathcal{O}(e^{-2kL})$ corrections. In the WTCM, $k e^{-kL} \sim 15$ TeV over a wide range of parameters, while the radion mass is given by the approximate relation [5]

$$m_\phi \approx \frac{k}{\bar{M}_P} (4 \text{ GeV}) \sim \left(\frac{35 \text{ TeV}}{\Lambda_\phi} \right) (4 \text{ GeV}). \quad (2)$$

Reliable gravity calculations and naturalness require that $k \lesssim M_5$, and hence we may expect $m_\phi \lesssim 4$ GeV, on general grounds, in our analysis. Typical values are $\Lambda_\phi \sim 100$ TeV and $m_\phi \sim 1$ GeV (with $kL \sim 30$).

We will next provide estimates of the rate for $b \rightarrow s\phi$ in the context of WTCM. There are two types of possible contributions to this process that we will consider in turn. First, there is a loop contribution, which is the

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analogue of the SM process with a light Higgs [9] instead of the radion. Secondly, there is a *tree-level* contribution that arises in 5D models that account for the SM fermion flavor, by localizing the quarks and leptons along the fifth dimension [10]. This effect is a consequence of the non-universality of the radion coupling to fermions with different 5D profiles and has recently been studied in Ref. [11]. Bulk localized fermions provide well-motivated flavor models, and the WTCM itself is based on this framework [5]. Hence, the effects of radion-coupling non-universalities are considered typical in this work.

Let us first look at the loop effect. The properties of the radion closely resemble those of the SM Higgs scalar h , and this allows us to adapt existing calculations for the process $B \rightarrow X_s h$ for our purposes. We will follow the results of Ref. [9] in our analysis. For B mesons, one can use the spectator quark approximation so that hadronic matrix elements cancel in the ratio $\text{Br}(B \rightarrow X_s \phi)/\text{Br}(B \rightarrow X_c e \bar{\nu}_e) \approx \text{Br}(b \rightarrow s \phi)/\text{Br}(b \rightarrow c e \bar{\nu}_e)$, and therefore (we assume IR-localized top quarks)

$$\frac{\text{Br}(B \rightarrow X_s \phi)}{\text{Br}(B \rightarrow X_c e \bar{\nu}_e)} = \frac{27\sqrt{2}}{64\pi^2} \frac{G_\phi m_b^2}{f(m_c/m_b)} \left| \frac{V_{st}^\dagger V_{tb}}{V_{cb}} \right|^2 \left(\frac{m_t}{m_b} \right)^4 \times \left(1 - \frac{m_\phi^2}{m_b^2} \right)^2, \quad (3)$$

where fermion masses and the CKM matrix elements are in obvious notation, $G_\phi \equiv (\sqrt{2}\Lambda_\phi^2)^{-1}$, and $f(m_c/m_b) \simeq 0.5$. Using $m_\phi \sim 1$ GeV and $\Lambda_\phi \sim 100$ TeV, we find

$$\text{Br}(B \rightarrow X_s \phi) \sim 10^{-5} \quad (\text{Loop}). \quad (4)$$

Next, we turn our attention to an additional contribution that can arise at tree-level, in warped models with 5D SM field content. A simultaneous resolution of the hierarchy and flavor puzzles of the SM can be achieved in these models, making them quite well-motivated. Quite generally, once the fermions are allowed to propagate in the 5D bulk, their interactions with massive bulk modes, such as the radion, are subject to non-universality. That is, once the fermion mass matrix is diagonalized, these bulk interactions, governed by the overlap of various 5D wavefunctions, are not necessarily diagonal and result in the appearance of FCNC's, that, in the case of the radion, can be parametrized as [11]

$$\mathcal{L}_{FV} = \frac{\phi}{\Lambda_\phi} (\bar{d}_L^i d_R^j a_{ij} \sqrt{m_i m_j} + \text{h.c.}), \quad (5)$$

where we have $i \neq j$. Here, a_{ij} is a function of the fermion zero-mode profiles [10], parametrized by the quantity $c \equiv M_\psi/k$; M_ψ is the 5D fermion mass. There are analogous interactions involving the up-type quark sector.

In the WTCM, the light fermion masses arise from 5D 4-fermion operators of the form¹

$(d_{ij}/\Lambda^3) (\bar{\Psi}_{Q_L^i} \Psi_{q_R^j}) (\bar{\Psi}_{Q_L^3} \Psi_{t_R}) + \text{h.c.}$ that lead at low energies to a Yukawa interaction with a scalar bound state $H \sim \bar{\Psi}_{Q_L^3} \Psi_{t_R}$. The effective 4D Yukawa matrices read $[N_c \bar{g}/(8\pi^2)] [d_{ij}/(\Lambda L)] [M_{KK}^2/\tilde{\Lambda}^2] f_{Q_L^i q_R^j Q_L t_R}$, where $\tilde{\Lambda} = \Lambda e^{-kL}$ is the warped down cutoff, $N_c = 3$ is the number of colors, $\bar{g} \approx 375/174$ is related to the dynamical quark mass generated by the condensation mechanism, and $f_{Q_L^i q_R^j Q_L t_R}$ is an overlap integral of four fermion wavefunctions, and is given in Section 3.3 of Ref. [5]. Since Q_L^3 and t_R are required to be near the IR brane to trigger the condensation (e.g. $c_{Q^3} \approx -2/3$, $c_{t_R} \approx -1/2$), the latter overlap integral is to a good approximation proportional to the product of the light fermion wavefunctions evaluated on the IR brane: $f_{Q_L^i q_R^j Q_L t_R} \approx (1/2) f_{c_{Q^i}}(L) f_{c_{q_R^j}}(L)$, where the zero-mode wavefunctions are given, for example, in Eq. (2) of Ref. [5]. It follows that with an anarchy assumption for the arbitrary coefficients d_{ij} , the flavor structure of the WTCM is similar to the widely studied RS models with fermions in the bulk. For instance, the CKM matrix elements are approximately given by $V_{ij}^{CKM} \sim f_{c_{Q^i}}(L)/f_{c_{Q^j}}(L)$ for $i < j$, thus fixing the $SU(2)$ doublet localization parameters, c_{Q^i} for $i = 1, 2$ [12]. The localization parameters for the RH fields can then be fixed so as to reproduce the fermion masses. For instance, an NDA estimate of the unknown coefficients of the 4-fermion interactions gives $d_{ij} \sim 24\pi^3/n$, where n is a measure of the number of fields appearing in loops. For $n = 10$, we find that with $c_{Q^1} = 0.65$, $c_{Q^2} = 0.58$, $c_{u^1} = 0.6$, $c_{u^2} = -0.4$, $c_{d^1} = 0.52$, $c_{d^2} = 0.45$, $c_{d^3} = 0.25$, the fermion masses and CKM mixing angles can be accommodated under the anarchy assumption. In this case, we find $a_{23} \approx a_{32} \approx 0.02$. For smaller n , the required localization parameters are very similar to the standard RS implementation of flavor (with all light fermions closer to the UV brane). Radion mediated FCNC's in such a case have been recently studied in Ref. [11], where a scan revealed typical values of $a_{ij} \sim 0.05$. We will take this as a representative value in the WTCM model.

Considering only the tree-level interaction (5), for the ϕbs vertex, we find

$$\frac{\text{Br}(B \rightarrow X_s \phi)}{\text{Br}(B \rightarrow X_c e \bar{\nu}_e)} = \frac{6\pi^2 a_{bs}^2 (m_s/m_b^3)(1 - m_\phi^2/m_b^2)^2}{\Lambda_\phi^2 G_F^2 |V_{cb}|^2 f(m_c/m_b)}, \quad (6)$$

where $a_{bs}^2 = |a_{23}|^2 + |a_{32}|^2$ and we have ignored terms of $\mathcal{O}(m_s/m_b)$ [in particular, we neglect a term proportional to $\text{Re}(a_{23} a_{32}^*)$]. Using as our reference values $m_\phi \sim 1$ GeV, $\Lambda_\phi \sim 100$ TeV, and $a_{ij} \sim 0.05$, we find

$$\text{Br}(B \rightarrow X_s \phi) \sim 10^{-2} \quad (\text{Tree}). \quad (7)$$

the relevant operators involve a linear combination of the RH top field and a new quark field with the same quantum numbers as t_R . However, this is not important in the present discussion.

¹ The implementation of a top seesaw in the WTCM implies that

Hence, we see that for typical values of WTCM parameters, the contribution from tree-level non-universality is larger by a factor of $\mathcal{O}(10^3)$, compared to the reference loop-level estimate obtained in Eq. (4). The loop effect is, in principle, less model dependent, since it does not rely on bulk fermion profiles and the resulting off-diagonal radion couplings. However, as noted before, fermion localization in warped models is well-motivated and provides a predictive and natural mechanism for explaining the SM flavor puzzle. Therefore, we adopt the point of view that, in the context of WTCM, the $b \rightarrow s\phi$ channel can typically have branching fractions of the order given in Eq. (7). Given the relatively large size of the branching fraction for this mode, we expect that it can provide distinct signatures that place bounds on the model and can potentially lead to its discovery, using B -decay data.

To go further in our phenomenological analysis, we must address the eventual decay of the radion. We will focus on $m_\phi \lesssim 4$ GeV, required for on-shell ϕ decays. However, given Eq. (2), this choice also coincides with the expected values of m_ϕ in the WTCM. Over the narrow range $3.7 \text{ GeV} \lesssim m_\phi \lesssim 4 \text{ GeV}$, ϕ decays mostly into $\tau^+\tau^-$ and $c\bar{c}$. For $m_\phi \lesssim 3.7 \text{ GeV}$, the most relevant decay modes of ϕ are into $\mu^+\mu^-$, $s\bar{s}$, and gg . Note that even though the gluon g is massless, the quantum trace anomaly of QCD allows for a non-negligible ϕgg coupling. In the perturbative QCD regime, which we roughly take to correspond to $m_\phi \gtrsim 1 \text{ GeV}$, using the results of Refs. [13, 14], we obtain

$$\Gamma_{\mu^+\mu^-} : \Gamma_{s\bar{s}} : \Gamma_{gg} \simeq x_\mu^2 m_\mu^2 : 3 x_s^2 m_s^2 : \left(\frac{b_4 \alpha_s}{2\pi} \right)^2 m_\phi^2, \quad (8)$$

where $m_s \simeq 104 \text{ MeV}$, $b_4 = 25/3$ is the QCD β -function for 4 flavors, and $\alpha_s/\pi \simeq 0.1$. Here $x = (c_L + c_R)_f$, where $f = \mu, s$, and $c_{L,R}$ parametrize the localization of f helicities in the extra dimension; quite typically we expect $x \approx 1$. We note that, in this work, brane kinetic terms are assumed to be generated at the quantum level and are hence ignored in radion couplings.

It then follows that, in the range $1 \text{ GeV} \lesssim m_\phi \lesssim 3.7 \text{ GeV}$, the radion almost exclusively decays into the gg final state. However, given the large suppression scale Λ_ϕ of ϕ couplings, we must first check that $\phi \rightarrow gg$ typically happens within the detector. Otherwise, the signal would resemble that of $b \rightarrow s\bar{t}$. To get an estimate, note that the lifetime τ_ϕ of the radion is roughly given by

$$\tau_\phi \sim \frac{32\pi^3 \Lambda_\phi^2}{b_4^2 \alpha_s^2 m_\phi^3} \sim 10^{-12} \text{ s} \left(\frac{\Lambda_\phi}{10^2 \text{ TeV}} \right)^2 \left(\frac{1 \text{ GeV}}{m_\phi} \right), \quad (9)$$

where $\mathcal{O}(1)$ factors have been ignored. The above lifetime τ_ϕ roughly corresponds to a displaced vertex of $\mathcal{O}(0.3) \text{ mm}$. Hence, we see that typical values of the parameters we are considering result in ϕ decays well within the detector. In addition, the displacement of the vertex is in a range that could be detected [15, 16] as a distinct signature for $\Lambda_\phi \gtrsim 100 \text{ TeV}$.

Therefore, in the range $1 \text{ GeV} \lesssim m_\phi \lesssim 3.7 \text{ GeV}$, the interaction in Eq. (5) results in a branching fraction

$$\text{Br}(b \rightarrow s\phi \rightarrow sgg) \sim 10^{-2} \quad (\text{WTCM}). \quad (10)$$

To extract a bound on the model parameters, we compare the above value to that expected from the SM [17]

$$\text{Br}(b \rightarrow sgg) \sim 10^{-3} \quad (\text{SM}). \quad (11)$$

Ignoring a weak dependence on m_ϕ , Eq. (6) then suggests

$$|a_{bs}|/\Lambda_\phi \lesssim (10^4 \text{ TeV})^{-1}. \quad (12)$$

Below 1 GeV, the radion decay enters a non-perturbative regime. Here, the dominant ϕ decay modes are expected to be $\pi^\pm, 0 \pi^\mp, 0$ (or $K^\pm, 0 K^\mp, 0$, if kinematically allowed), since we expect that these hadronic decays are controlled by the coupling to gluons [9]. Given the inherent uncertainty in the size of this effect, due to unknown chiral Lagrangian coefficients [9], we will not present any numerical bounds here. In any event, the dominant signal for $2m_\pi \lesssim m_\phi \lesssim 1 \text{ GeV}$ is expected to be of the form $X_s \pi \pi$, with displaced $\pi\pi$ vertices.

In the narrow region $2m_\mu < m_\phi < 2m_\pi$, the dominant decay modes of the radion are $\mu^+\mu^-$ and $\gamma\gamma$. Here, unlike the gluons, we find that the final state photons mostly couple through their bulk overlap with the radion. To see this, note that the QED β -function $b_{\text{QED}} = 16/3$ below the τ -charm threshold [14], and due to the smallness of $\alpha \simeq 1/137$, the quantum effect is negligible. We then expect that the radion partial widths will be given by

$$\Gamma_{\mu^+\mu^-} : \Gamma_{\gamma\gamma} \simeq m_\mu^2 : m_\phi^2/(kL)^2. \quad (13)$$

Since $m_\phi \approx 2m_\mu$ over this mass interval and $kL \simeq 30$ in the WTCM, the di-muon final state dominates in this narrow range of masses. Experimental data [18] give

$$\text{Br}(B \rightarrow s\mu^+\mu^-) = (4.3 \pm 1.2) \times 10^{-6} \quad (\text{Data}). \quad (14)$$

Assuming that the SM is only perturbed by the new physics, one can then demand that the size of the effect from WTCM not be larger than the error on this measurement. This then yields

$$|a_{bs}|/\Lambda_\phi \lesssim (2 \times 10^5 \text{ TeV})^{-1}, \quad (15)$$

which is a severe bound on the model parameters. Note that in this regime, if we keep $|a_{bs}| \sim 0.05$, the effective radion coupling is pushed to values of order 10^4 TeV . For such a high value of the radion coupling scale, we expect $\tau_\phi \sim 16\pi\Lambda_\phi^2/(m_\mu^2 m_\phi) \sim 10^{-9} \text{ s}$, which means that the decay starts to occur outside the detector. However, this regime is beyond the expected WTCM value $\Lambda_\phi \lesssim 10^3 \text{ TeV}$ from Eq. (2). On the other hand, keeping the radion coupling near the $\mathcal{O}(10^3) \text{ TeV}$ range suggests that the flavor parameters may need to be tuned.

The dominant decay mode, over the interval $15 \text{ MeV} \lesssim m_\phi \lesssim 2m_\mu$, is the di-photon mode. This can be deduced

from Eq. (13) with $m_\mu \rightarrow m_e$, with $kL \lesssim 30$. Note that near the upper end of this regime, we expect $\Lambda_\phi \sim 10^3$ TeV and $kL \sim 30$. Given

$$\tau_\phi \sim 16\pi(kL)^2\Lambda_\phi^2/m_\phi^3 \quad (16)$$

we find that $\tau_\phi \gtrsim 10^{-6}$ s and this mass interval will be characterized by a $b \rightarrow s\bar{H}$ signal. We then find that the branching fraction (6) for this mode will be below 10^{-4} for $|a_{bs}| \sim 0.05$. Given that the typical experimental bounds on the branching fraction for $B \rightarrow X_s \nu\bar{\nu}$ are at this level [18], we do not expect a severe bound in this regime from B -decay data.

So far, we have only addressed possible laboratory bounds on this scenario. However, for a sufficiently light radion, astrophysical bounds could become important. In particular, for $m_\phi \lesssim 30$ MeV, radions are light compared to the supernova core temperature of $\mathcal{O}(30)$ MeV and can provide an energy loss mechanism for the core. For this to happen, the decay length of the radion, set by $\phi \rightarrow \gamma\gamma$, must exceed $\mathcal{O}(100)$ km, the typical size of a type II supernova core. We find this to be the case, in the above mass range.

Thus, for $m_\phi \lesssim 30$ MeV, astrophysical over cooling bounds from SN 1987A can be important. For example, light radions can be produced in bremsstrahlung processes of the type $NN \rightarrow NN\phi$, where N denotes a nucleon. Given a radion coupling of order m_N/Λ_ϕ , SN 1987 A bounds require $\Lambda_\phi \gtrsim 10^6$ TeV, where we have adapted similar bounds on axion emission [18] to the case of scalar emission [19]. This corresponds to $m_\phi \sim 100$ keV. The onset of over cooling constraints cor-

responds to $m_\phi \lesssim 30$ MeV and roughly $\Lambda_\phi \gtrsim 10^4$ TeV. We thus conclude that the approximate range $0.1 \text{ MeV} \lesssim m_\phi \lesssim 30 \text{ MeV}$, is disfavored by astrophysical data on SN 1987 A. From a theoretic point of view, one could also argue that very small values of k/\bar{M}_P reintroduce the kind of hierarchy that the WTCM is meant to address and therefore such low radion masses are not well-motivated in this framework [see Eq. (2)].

In summary, we have considered the light radion phenomenology of the recently proposed WTCM in the context of B physics data. We found that the tree-level vertex $bs\phi$, from radion-fermion coupling non-universality, provides the dominant contribution to $b \rightarrow s\phi$. Over much of the natural parameter space of the model, the radion will decay into light hadrons or muons, with a measurable displaced vertex. Various bounds on the model parameters, based on theoretical expectations within the SM and current experimental data, were obtained. A more detailed analysis of the current and future data, using the vertex displacement information can either enhance the bounds we have found or lead to possible discovery of the light radion signal. We also briefly considered possible astrophysical bounds that may apply for very light radions.

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